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# Possible explanation of the electron positron anomaly at 17 MeV in $^8\text{Be}$ transitions through a light pseudoscalar

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**ABSTRACT:** We estimate the values of Yukawa couplings of a light pseudoscalar  $A$  with a mass of about 17 MeV, which would explain the  $^8\text{Be}$  anomaly observed in the Atomki pair spectrometer experiment. The resulting couplings of  $A$  to up and down type quarks are about 0.3 times the coupling of the standard Higgs boson. Then constraints from  $K$  and  $B$  decays require that loop contributions to flavour changing vertices cancel at least at the 10% level. Constraints from beam dump experiments require the coupling of  $A$  to electrons to be larger than about 4 times the coupling of the standard Higgs boson, leading to a short enough  $A$  life time consistent with an explanation of the anomaly.

**KEYWORDS:** Phenomenological Models, Supersymmetry Phenomenology

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## 1 Introduction

The Atomki pair spectrometer experiment [1] has searched for electron-positron internal pair creation in the decay of excited  ${}^8\text{Be}$  nuclei. The  ${}^8\text{Be}$  excitations were produced with help of a beam of protons directed on a  ${}^7\text{Li}$  target and the different  ${}^8\text{Be}$  excitations could be separated by tuning the energy of the incoming protons.

An anomaly has been observed in the decay of  ${}^8\text{Be}^*$  with spin-parity  $J^P = 1^+$  into the ground state  ${}^8\text{Be}$  with spin-parity  $0^+$  (both with isospin  $T = 0$ ), where  ${}^8\text{Be}^*$  has an excitation energy of 18.15 MeV. Both distributions of the opening angle  $\theta$  of the electron-positron pair and the invariant mass of the electron-positron pair showed an excess consistent with an intermediate boson  $X$  being produced in the decay of  ${}^8\text{Be}^*$ , with  $X$  decaying into an electron-positron pair. The best fit to the mass  $M_X$  of  $X$  is [1]

$$M_X = 16.7 \pm 0.35 \text{ (stat)} \pm 0.5 \text{ (sys)} \text{ MeV} \quad (1.1)$$

whereas the best fit to the branching fraction  ${}^8\text{Be}^* \rightarrow {}^8\text{Be} + (X \rightarrow e^+e^-)$  relative to the branching fraction  ${}^8\text{Be}^* \rightarrow {}^8\text{Be} + \gamma$  is given by

$$\frac{Br({}^8\text{Be}^* \rightarrow X + {}^8\text{Be}) \times Br(X \rightarrow e^+e^-)}{Br({}^8\text{Be}^* \rightarrow \gamma + {}^8\text{Be})} = 5.8 \times 10^{-6} . \quad (1.2)$$

These values correspond to a statistical significance of the excess of  $6.8\sigma$  [1].

In the case of the excitation  ${}^8\text{Be}^{*'} with spin-parity  $1^+$  (but isospin  $T = 1$ ) and an excitation energy of 17.64 MeV, no excess was observed. The simplest explanation is that this decay is kinematically suppressed; this kinematical suppression is the stronger the heavier the intermediate boson  $X$  would be. This motivates a value of  $M_X$  somewhat above the best fit value in (1.1) (which may lead to a somewhat smaller statistical significance and smaller best fit to the relative branching fraction).$

In [2, 3] an explanation for the observed excess was given in the form of models featuring a new vector boson  $Z'_\mu$  with a mass  $M_{Z'}$  of about 17 MeV, with vector-like couplings to

quarks and leptons. Constraints on such a new vector boson, notably from searches for  $\pi^0 \rightarrow Z' + \gamma$  by the NA48/2 experiment [4], require that the couplings of  $Z'_\mu$  to up and down quarks are “protophobic”, i.e., that the charges  $e\varepsilon_u$  and  $e\varepsilon_d$  of up and down quarks — written as multiples of the positron charge  $e$  — satisfy  $2\varepsilon_u + \varepsilon_d \lesssim 10^{-3}$  [2, 3]. Subsequently, further studies of such models have been performed in [5–8].

Given the quantum numbers of the  ${}^8\text{Be}^*$  and  ${}^8\text{Be}$  states, the boson  $X$  can also be a pseudoscalar  $A$  with a mass  $M_A$  of about 17 MeV. In [2, 3] this possibility is dismissed quite rapidly. The argument is that, for such an axion-like pseudoscalars  $A$ , fermion loops generate couplings of the form  $g_{A\gamma\gamma} A F^{\mu\nu}(\gamma) \tilde{F}_{\mu\nu}(\gamma)$  which are strongly constrained by axion searches. However, light pseudoscalars in this mass range with tree level Yukawa couplings to electrons decay dominantly into electron-positron pairs, unless Yukawa couplings to other charged fermions  $f$  with mass  $m_f$  are much larger than  $m_f/m_e$  compensating  $g_{A\gamma\gamma} \approx 1/(8\pi m_f)$ .

It is the purpose of the present paper to study the required couplings of a pseudoscalar  $A$  with a mass of about 17 MeV in order to explain the  ${}^8\text{Be}$  anomaly observed in [1], and to verify under which conditions these couplings satisfy existing constraints. We have in mind a pseudoscalar  $A$  originating from extended Higgs sectors of the Standard Model (SM) including, e.g., two Higgs doublets of type II and a singlet as in the Next-to-Minimal Supersymmetric SM (NMSSM) [9], where  $A$  could be very light in Peccei-Quinn or  $R$ -symmetry limits [9]. We find however that (singlet extended) two Higgs doublet models of type II have difficulties to explain the anomaly, but more general models are possible under the condition that the various loop contributions to the flavour changing vertex  $A - s - d$  cancel at least at the 10% level.

A major task is to express the coupling of such a pseudoscalar to  ${}^8\text{Be}^*$  and  ${}^8\text{Be}$  states in terms of the couplings of  $A$  to up and down quarks. Required is actually the ratio of branching fractions

$$\frac{Br({}^8\text{Be}^* \rightarrow A + {}^8\text{Be}) \times Br(A \rightarrow e^+ e^-)}{Br({}^8\text{Be}^* \rightarrow \gamma + {}^8\text{Be})} \quad (1.3)$$

which is given in (1.2). In the case of the  $Z'$  considered in [2, 3], use is made of the fact that both  $Z'$  and photons couple via conserved currents to quarks, an argument which is not useful here. Furthermore, [2, 3] argue that both  $Z'_\mu$  and photons couple via conserved currents to nucleons, and that — at least in the isospin conserving limit considered in [2] — matrix elements of conserved currents cancel in the calculation of the ratio of decay widths up to the modifications of the couplings. (The possible impact of isospin violating effects is analysed in [3].)

The calculation of the coupling of a pseudoscalar  $A$  to  ${}^8\text{Be}^*$  and  ${}^8\text{Be}$  states has to proceed in two steps. Firstly, the couplings of  $A$  to nucleons have to be obtained: these are proportional to the nucleon quark spin components  $\Delta q$ , and have been studied in the context of direct detection of dark matter via the exchange of pseudoscalars, e.g., in [10, 11]. Secondly, the  ${}^8\text{Be}^*$  and  ${}^8\text{Be}$  nuclei have to be described in terms of nucleons with definite spin, angular momentum and total momentum. To this end we employ wave functions from the simple unperturbed nuclear shell model. We are aware of the fact that this approach is somewhat simplistic: it neglects proton-neutron pairing effects,  $\alpha - \alpha$  substructures of the

${}^8\text{Be}$  states and, in particular, possible mixing with the nearby  ${}^8\text{Be}^{*'}$  state induced by isospin breaking. Effects of the latter have been discussed in [3], and could be sizeable. For consistency, we have to employ the same approach for the decay widths  $\Gamma({}^8\text{Be}^* \rightarrow \gamma + {}^8\text{Be})$  and  $\Gamma({}^8\text{Be}^* \rightarrow A + {}^8\text{Be})$ . One may hope that the inaccuracies of the nuclear shell model wave functions cancel to some extent in the calculations of the ratio of decay widths, but we will return to this issue later on. In any case some theoretical error has certainly to be taken into account, and a further refinement of the present calculation of this ratio would be desirable.

The plan of the paper is as follows. In section 2 we consider the couplings of a pseudoscalar to nucleons while in section 3 we compute and compare the relevant matrix elements for  $\gamma$  and pseudoscalar emission in the nuclear shell model. In this section we also find the conditions on the pseudoscalar Yukawa couplings to quarks and leptons which are necessary in order to explain the anomaly. Section 4 is devoted to other experimental constraints on these couplings. Finally, a summary and some conclusions are presented in section 5.

## 2 Couplings of a pseudoscalar to nucleons

Subsequently we define reduced couplings  $\xi_q$  of a pseudoscalar  $A$  to quarks in terms of

$$\mathcal{L}_{Aqq} = \xi_q \frac{m_q}{v} A \bar{q} i \gamma_5 q \quad (2.1)$$

with  $v \sim 246 \text{ GeV}$ . As in [10] we define a pseudoscalar-nucleon coupling  $h_N$  (with  $N = p, n$  for protons and neutrons, respectively) by

$$h_N = \frac{1}{v} \sum_q \langle N | \xi_q m_q \bar{q} i \gamma_5 q | N \rangle . \quad (2.2)$$

From [10] (see also [11]) one finds

$$h_N = \frac{m_N}{v} \sum_{q=u,d,s} \Delta_q^{(N)} \left( \xi_q - \sum_{q'=u,\dots,t} \xi_{q'} \frac{\bar{m}}{m_q} \right) , \quad (2.3)$$

where  $\Delta_q^{(N)}$  are the quark spin components of the nucleon  $N$ , and  $\bar{m} = \frac{1}{m_u^{-1} + m_d^{-1} + m_s^{-1}} \sim \frac{m_u m_d}{m_u + m_d}$ . In addition, we assume [10]  $m_d \sim 2m_u \sim 2 \times 2.5 \text{ MeV}$  and

$$\xi_u = \xi_c = \xi_t, \quad \xi_d = \xi_s = \xi_b . \quad (2.4)$$

Neglecting  $\frac{m_{u,d}}{m_{s,c,b,t}}$  one obtains

$$h_N = \frac{m_N}{v} \left( \Delta_u^{(N)} (-\xi_u - 2\xi_d) + \Delta_d^{(N)} (-\xi_u) + \Delta_s^{(N)} \xi_d \right) . \quad (2.5)$$

For  $\Delta_q$  we use the values given in table II in [10] using  $g_A^8 = 0.46$  and  $g_A^0 = 0.37$ :

$$\Delta_u^{(p)} = 0.84, \quad \Delta_d^{(p)} = -0.44, \quad \Delta_s^{(p)} = -0.03, \quad \Delta_u^{(n)} = -0.44, \quad \Delta_d^{(n)} = 0.84, \quad \Delta_s^{(n)} = -0.03 . \quad (2.6)$$

This gives

$$h_p = \frac{m_p}{v} (-0.40\xi_u - 1.71\xi_d), \quad h_n = \frac{m_n}{v} (-0.40\xi_u + 0.85\xi_d) . \quad (2.7)$$

For the average  $\bar{h}_N^2 \equiv \frac{(h_p+h_n)^2}{4}$ , required for the  ${}^8\text{Be}^*$  decays, one obtains (with  $m_n \sim m_p$ )

$$\bar{h}_N^2 = \frac{m_p^2}{v^2} f(\xi_u, \xi_d), \quad f(\xi_u, \xi_d) = (0.16\xi_u^2 + 0.35\xi_u\xi_d + 0.19\xi_d^2) . \quad (2.8)$$

For  $\xi_u = \xi_d \equiv \xi$  one has  $f(\xi, \xi) \sim 0.7\xi^2$ .

### 3 Nuclear shell model and emission matrix elements

The  ${}^8\text{Be}$  ground state with  $J^P = 0^+$  and the  ${}^8\text{Be}^*$  excited state with  $J^P = 1^+$  can be described in terms of the lowest two shells of the nuclear shell model: the lowest 1s ( $L = 0$ ) shell is fully occupied by two nucleons with spin  $S_z = \pm 1/2$  (two out of the four protons and two out of the four neutrons); in the next 1p ( $L = 1$ ) shell there is, a priori, space for six nucleons with angular momentum  $L_z = -1, 0, +1$  and  $S_z = \pm 1/2$ , respectively. However, the spin-orbit interaction proportional to  $-\langle \vec{L} \cdot \vec{S} \rangle$  splits the 1p level into two levels with total angular momentum  $J = 3/2$  (four possible states  $1p_{3/2}$ ) and  $J = 1/2$  (two possible states  $1p_{1/2}$ ) where the  $J = 3/2$  level is lower. In the  ${}^8\text{Be}$  ground state two out of the four  $1p_{3/2}$  states are occupied by protons/neutrons respectively, and the angular momenta can be combined pairwise to form a nucleus with  $J^P = 0^+$ .

If one of the two states in the lower  $1p_{3/2}$  level is lifted into the previously empty  $1p_{1/2}$  level it would form with its remaining partner in the  $1p_{3/2}$  level a  $J^P = 1^+$  state which gives, together with the remaining  $J^P = 0^+$  nucleons, a  $J^P = 1^+$  state consistent with the quantum numbers of  ${}^8\text{Be}^*$ . Its excitation energy of 18.15 MeV is consistent with — following [12] perhaps slightly larger than — the expectations from nuclear spin-orbit splitting. During the transition from  ${}^8\text{Be}^*$  to  ${}^8\text{Be}$  a photon or — as considered here — a pseudoscalar can be emitted from a single nucleon falling from a  $1p_{1/2}$  state into the lower  $1p_{3/2}$  state. The photon emission is of the M1 type.

The next task is to construct the interaction Hamiltonian for both M1 photon and pseudoscalar emissions from single  $1p_{1/2}$  nucleon states; finally we need the ratios of both decay rates which should be compared — together with the  $A \rightarrow e^+e^-$  branching fraction — to  $5.8 \times 10^{-6}$  (1.2) as estimated for the signal in [1].

In order to treat the photon and pseudoscalar emissions at the same level we construct first the non-relativistic interaction Hamiltonian from the relativistic Dirac equation for single nucleons  $N = p, n$ . After adding a coupling  $h_N$  to a pseudoscalar  $A$  and an anomalous magnetic moment  $\sim (g - 2)_N$  to the Lagrangian, the Dirac equation including the covariant  $U(1)_{\text{em}}$  derivative with a photon  $A^\mu = (\phi, A^i)$  can be written as (isolating the time derivative)

$$i\hbar\gamma^0\partial_t\psi_N = (\gamma^i(p^i - q_N A^i) + \gamma^0 q_N \phi + \frac{(q \cdot (g - 2))_N}{8m_N} \sigma_{\mu\nu} F^{\mu\nu} + m_N + i h_N A \gamma^5) \psi_N + \dots \quad (3.1)$$

where the dots describe the potential (including spin-orbit terms etc.) for single nucleons generated by the seven remaining nucleons of  ${}^8\text{Be}$ .

Decomposing  $\psi_N = \begin{pmatrix} \tilde{\varphi}_N \\ \tilde{\chi}_N \end{pmatrix}$ ,  $\tilde{\varphi}_N = e^{im_N t} \varphi_N$ ,  $\tilde{\chi}_N$  can be eliminated in the non-relativistic limit in an expansion in  $1/m_N$ . To lowest order in the couplings  $e$ ,  $h_N$  the remaining Schrödinger equation for  $\varphi_N$  contains an interaction Hamiltonian of the form

$$H_{\text{int}} = -q_N \vec{r} \cdot \vec{E} - \frac{1}{2m_N} \left( q_N \vec{B} \cdot \vec{L} + (q \cdot g)_N \vec{B} \cdot \vec{S} + 2h_N (i\vec{\nabla} A) \cdot \vec{S} \right) \quad (3.2)$$

where the first term is irrelevant for M1 transitions, and  $\vec{S} = \frac{1}{2}\vec{\sigma}$ . In (3.2) the  $g$ -factor  $(q \cdot g)_N$  includes the anomalous magnetic moment  $\sim (g - 2)$ : for protons one has to use  $q_p = e$ ,  $g_p = 5.6$ , for neutrons  $q_n = 0$  in the first terms, but  $(q \cdot g)_n = -3.8e$ . The coupling of the pseudoscalar  $A$  is as expected:  $\vec{\nabla} A$  indicates that  $A$  can be emitted only as a  $p$ -wave, and couples to the spin.

Next one has to evaluate the matrix elements of  $H_{\text{int}}$  between the states  $\langle J' = 3/2, m_{j'} |$  and  $| J = 1/2, m_j \rangle$ ; the decay rates are proportional to

$$\sum_{m_{j'}} |\langle 3/2, m_{j'} | H_{\text{int}} | 1/2, m_j \rangle|^2 \quad (3.3)$$

where one has to average over  $m_j = \pm 1/2$ . From the different terms in the decay rates one can estimate the ratio between photon and pseudoscalar emission.

Let us emit the photon with momentum  $\vec{p}_\gamma$  and the pseudoscalar with momentum  $\vec{p}_A$  in the  $z$  direction, leading to  $|B_x|^2 = |B_y|^2 = \vec{p}_\gamma^2 |A_\mu|^2$ . Then one finds for (3.3) (still for a given nucleon  $N$ )

$$\begin{aligned} \frac{1}{4m_N^2} \sum_{m_{j'}} \left[ 2B_x^2 |q_N \langle 3/2, m_{j'} | L_x | 1/2, m_j \rangle + (q \cdot g)_N \langle 3/2, m_{j'} | S_x | 1/2, m_j \rangle|^2 \right. \\ \left. + 4h_N^2 \vec{p}_A^2 A^2 |\langle 3/2, m_{j'} | S_z | 1/2, m_j \rangle|^2 \right], \end{aligned} \quad (3.4)$$

and, finally, after evaluating the matrix elements of  $L_x$ ,  $S_x$  and  $S_z$ ,

$$\frac{1}{2} \sum_{m_{j'}, m_j} |\langle 3/2, m_{j'} | H_{\text{int}} | 1/2, m_j \rangle|^2 = \frac{1}{9m_N^2} \cdot \left[ \vec{p}_\gamma^2 |A_\mu|^2 (q_N - (q \cdot g)_N)^2 + 2h_N^2 \vec{p}_A^2 A^2 \right] \quad (3.5)$$

which, for isospin singlet nuclei, has to be averaged over the nucleon states  $N = p$  and  $N = n$  (including interference terms). The two terms  $\sim |A_\mu|^2$  and  $\sim A^2$  on the right hand side of (3.5) correspond to the emission of the photon  $\gamma$  and pseudoscalar  $A$ , respectively. Using the expressions given below eq. (3.2), the average of the coefficient  $(q_N - (q \cdot g)_N)^2$  becomes

$$\frac{1}{4} (q_p - (q \cdot g)_p + q_n - (q \cdot g)_n)^2 \simeq 0.16 e^2. \quad (3.6)$$

The average for pseudoscalar couplings  $2\bar{h}_N^2$  is from (2.8)

$$2\bar{h}_N^2 = \frac{2m_p^2}{v^2} f(\xi_u, \xi_d) = 2.92 \times 10^{-5} f(\xi_u, \xi_d). \quad (3.7)$$

The decay rates also depend on powers of the photon/pseudoscalar momenta which originate from the phase space and normalization of the plane waves  $A_\mu$  and  $A$ ; the final dependence on the momenta is  $\sim |\vec{p}|^3$  in both cases. For the ratio of the decay rates one obtains then

$$\frac{Br(^8Be^* \rightarrow ^8Be + A)}{Br(^8Be^* \rightarrow ^8Be + \gamma)} = \frac{2.92 \times 10^{-5} f(\xi_u, \xi_d)}{0.16 e^2} \frac{|\vec{p}_A|^3}{|\vec{p}_\gamma|^3} = 2 \times 10^{-3} f(\xi_u, \xi_d) \frac{|\vec{p}_A|^3}{|\vec{p}_\gamma|^3}, \quad (3.8)$$

where  $e^2 \simeq 0.091$  was used. Assuming a  $Br(A \rightarrow e^+e^-) \sim 1$  (see below), this expression should give

$$\frac{Br(^8Be^* \rightarrow ^8Be + A)}{Br(^8Be^* \rightarrow ^8Be + \gamma)} \approx 5.8 \times 10^{-6}. \quad (3.9)$$

The ratio of momenta depends on  $M_A$ . Taking  $M_A = 17$  MeV leads to

$$\frac{|\vec{p}_A|^3}{|\vec{p}_\gamma|^3} \sim 0.045. \quad (3.10)$$

From the three previous equations one obtains

$$f(\xi_u, \xi_d) \stackrel{!}{\approx} 0.062. \quad (3.11)$$

Approximating  $f(\xi_u, \xi_d)$  by  $f(\xi_u, \xi_d) \sim 0.175 (\xi_u + \xi_d)^2$  gives

$$\xi_u + \xi_d \stackrel{!}{\approx} 0.6 \quad (3.12)$$

or, for  $\xi_u = \xi_d \equiv \xi$ ,  $\xi \stackrel{!}{\approx} 0.3$ .

One should keep in mind, however, that this result depends on the use of the nuclear shell model wave functions with definite isospin  $T = 0$ . In particular, the coefficient 0.16 on the right hand side of (3.6) originates from substantial cancellations in the case of isoscalar  $M1$  transition strengths, a phenomenon underlined before in [3]. If this coefficient turns out to be larger due to a  $T = 1$  component in the  $^8Be^*$  wave function, the resulting value for  $f(\xi_u, \xi_d)$  in (3.11) increases as well. Of course, the expression for  $f(\xi_u, \xi_d)$  given in (2.8) would have to be corrected as well in this case, but here no strong cancellations occur in general. Hence the theoretical uncertainty to associate to the result (3.11) or (3.12) points towards rather larger values for  $\xi_u$  and/or  $\xi_d$  required to fit the anomaly observed in the Atomki pair spectrometer experiment.

We close this section with a consideration of the  $A$  width and decay length. If  $A$  has Yukawa couplings to quarks and leptons which are proportional to the Yukawa couplings of the SM Higgs boson rescaled by generation independent factors  $\xi_d \approx \xi_u \approx \xi_e$  (or  $\xi_u \ll \xi_d$ ), and the Yukawa couplings to BSM fermions are not much larger than the electric charge  $e$ ,  $A$  has a branching fraction of about 99% into  $e^+e^-$  and only about 1% into  $\gamma\gamma$ . Its total width is then dominated by  $A \rightarrow e^+e^-$  and given by

$$\Gamma(A) = \xi_e^2 \frac{m_e^2}{8\pi v^2} M_A = \xi_e^2 \cdot 2.9 \times 10^{-15} \text{ GeV} \quad (3.13)$$

for  $M_A = 17 \text{ MeV}$ . Its decay length is

$$l_A = \frac{p_A}{M_A \Gamma(A)} . \quad (3.14)$$

For the decay  ${}^8\text{Be}^* \rightarrow {}^8\text{Be} + A$  with  $M({}^8\text{Be}^*) - M({}^8\text{Be}) = 18.15 \text{ MeV}$  we obtain

$$l_A \sim \frac{1}{\xi_e^2} \cdot 2.5 \text{ cm} . \quad (3.15)$$

(For  $M_A = 17.9 \text{ MeV}$ ,  $2\sigma$  above the central value in (1.1), we obtain  $l_A \sim \frac{1}{\xi_e^2} \cdot 1.1 \text{ cm}$ .) In order to explain the observed anomaly in the Atomki pair spectrometer experiment [1],  $l_A$  should then not be much larger than 1 cm leading to

$$\xi_e \stackrel{!}{\gtrsim} 1 , \quad (3.16)$$

depending somewhat on the precise value of  $M_A$ .

## 4 Experimental constraints

Light pseudoscalars are subject to constraints from searches for axions or axion-like particles. For recent summaries of constraints relevant for light pseudoscalars decaying dominantly into  $e^+e^-$  see [11, 13–16]. However, since we allow for different Yukawa type couplings rescaled by  $\xi_u$ ,  $\xi_d$  and  $\xi_e$  with respect to SM Higgs couplings, at least some experimental constraints studied therein have to be reconsidered. We note that constraints from  $\pi^0 \rightarrow \gamma + X$  from the NA48/2 experiment, which play a major rôle for the  $Z'$  scenario [2, 3], do not apply here since the decay  $\pi^0 \rightarrow \gamma + A$  would violate parity. Furthermore, a light pseudoscalar cannot improve the discrepancy between the measured and the SM value of the anomalous magnetic moment of the muon since its contribution has the wrong sign (but is smaller in absolute value than the present discrepancy).

A first class of constraints on such pseudoscalars originates from flavour violating meson decays, analysed recently in [11]. For  $M_A \sim 17 \text{ MeV}$  and the range of couplings relevant here these are the decays  $K^+ \rightarrow \pi^+ + X$  (constrained by the  $K_{\mu 2}$  experiment [17]),  $K^+ \rightarrow \pi^+ + \textit{invisible}$  as measured by the experiments E787 [18] and BNL-E949 [19],  $B_s \rightarrow \mu^+ \mu^-$  (measured by the LHCb collaboration [20] and the CMS collaboration [21], see [22] for a LHCb/CMS combination), and  $B^0 \rightarrow K_S^0 + \textit{invisible}$  measured by CLEO [23].

Concerning  $K^+ \rightarrow \pi^+ + X$ , [17] searched for an anomalous line corresponding to  $\pi^+$  in the  $K_{\mu 2}$  experiment, which would appear for  $K^+ \rightarrow \pi^+ + A$  decays independently of subsequent  $A$  decays. This process depends on a loop-induced  $A - s - d$  vertex (with  $W$  bosons and up-type quarks in the loop, to be supplemented at least by  $H^\pm$  bosons in consistent multi-Higgs extensions of the SM) which depends, in turn, on the couplings of  $A$  to down and up type quarks (and to  $W^\pm H^\mp$ ).

Constraints from figure 2 in [17] have been applied to a light pseudoscalar in the NMSSM in [13]. Here squark/chargino loops are considered, which are dominant for large  $\tan \beta$  ( $\xi_d \gg \xi_u$ ) [24]. The resulting bound on  $C_{\text{Aff}}$  in [13] can be translated into  $\xi_d = C_{\text{Aff}}$ , which for  $M_A \sim 17 \text{ MeV}$  is

$$\xi_d \lesssim 2 \times 10^{-2} . \quad (4.1)$$



An even stronger bound has been derived in [11] in terms of  $g_Y$ , a common factor rescaling the Higgs-like Yukawa couplings of  $A$ . Note that  $\xi_u = \xi_d \equiv \xi$  corresponds to  $g_Y = \xi/\sqrt{2}$  in [11]. These authors find that  $g_Y \gtrsim 5 \times 10^{-3}$  or  $\xi \gtrsim 7.1 \times 10^{-3}$  is ruled out from [17]. However, the calculation of the loop-induced  $A - s - d$  vertex, relevant for  $K^+ \rightarrow \pi^+ + A$ , was performed in [11] without a charged Higgs boson in the loops leading to Ultra-Violet (UV) divergencies  $\sim \ln^2(\Lambda/m_{\text{top}})$ , a factor assumed to be of  $\mathcal{O}(10)$ . As discussed in [11], the divergencies are cancelled in UV complete models featuring a light pseudoscalar and in which the combined contributions to the  $A - s - d$  vertex can potentially be much smaller.

An example is provided by the similar process  $B \rightarrow K + A$  depending on the loop induced  $A - b - s$  vertex, studied in models of the two-Higgs-doublet (+ singlet) type in [25–27]. As it can be seen in [27] the partial width can vanish for appropriate choices of parameters (for  $M_{H^\pm} \sim 600 \text{ GeV}$  in two-Higgs-doublet models) due to cancellations in the loop functions. Up to different quark masses, the same loop functions appear in contributions to the  $A - s - d$  vertex. Also within supersymmetric extensions of the SM the a priori larger loop contributions to the  $A - s - d$  vertex [24] can cancel for, e.g., appropriate values of  $A_{\text{top}}$  and squark masses within the NMSSM [28]. We estimate that tunings at the 10% level within two-Higgs-doublet (+ singlet) models, but at most at the 1% level within supersymmetric extensions of the SM would be necessary in order to circumvent the upper bounds on  $\xi_d$  from  $K^+ \rightarrow \pi^+ + A$ . Albeit not elegant, the possibilities of such cancellations provide a go-theorem allowing for a light pseudoscalar to circumvent constraints from flavour changing processes in general.

Constraints from searches for  $K^+ \rightarrow \pi^+ + \text{invisible}$  from E787 and BNL-E949 [18, 19] apply only if  $A$  decays outside the detectors, i.e., if  $\xi_e$  is small enough. According to [13], identifying now  $C_{\text{Aff}}$  in [13] with  $\xi_e$ , this is not the case for  $\xi_e \gtrsim 0.3$ .

According to [11], the constraints from  $B_s \rightarrow \mu^+ \mu^-$  (through an off-shell  $A$ ) rule out  $g_Y \gtrsim 0.5$  or  $\xi \gtrsim 0.7$  which is weaker than the constraint (4.1) from  $K^+ \rightarrow \pi^+ + A$ . Again, the loop contributions to the  $A - s - b$  vertex considered in [11] are incomplete within a UV complete extension of the Higgs sector, and could again be cancelled by additional beyond-the-SM contributions as in the case of the  $A - s - d$  vertex.

The constraints from  $B^0 \rightarrow K_S^0 + \text{invisible}$  measured by CLEO [23] apply only if the pseudoscalar  $A$  produced in  $B^0 \rightarrow K_S^0 + A$  decays outside the detector. Accordingly these constraints depend both on the  $Br(B^0 \rightarrow K_S^0 + A)$ , hence on the  $A - s - b$  vertex or on  $\xi_u, \xi_d$ , and on the  $A$  decay length which depends on  $\xi_e$ . These quantities are identified in [11] where a limit  $g_Y \gtrsim 5$  or  $\xi \gtrsim 3.5$  satisfies the constraints, since then the  $A$  decay length becomes short enough despite the large production rate. Using this constraint only for  $\xi_e$  is conservative, if  $\xi_u, \xi_d < \xi_e$  is assumed.

Finally,  $\xi_e \gtrsim 3.5$  satisfies also bounds on  $A$  production in radiative  $\Upsilon$  decays  $\Upsilon \rightarrow \gamma + \text{invisible}$  interpreted as  $\Upsilon \rightarrow \gamma + A$  from CLEO [29] and BaBar [30], which apply only if  $A$  decays outside the detectors. For  $M_A \sim 17 \text{ MeV}$ , following [13], this is not the case for  $\xi_e \gtrsim 1.5$ .

A second class of constraints on light pseudoscalars originates from beam dump experiments, which we discuss in turn. First, an electron beam dump on lead experiment was conducted in Orsay [31] with the aim to search for light scalar or pseudoscalar Higgs

bosons in the decay into  $e^+e^-$ , produced via radiation off electrons. Correspondingly the resulting constraint applies to  $\xi_e$  only. According to [31] life times  $\tau_A$  in the range  $5 \cdot 10^{-12} \text{ s} \lesssim \tau_A \lesssim 2 \cdot 10^{-9} \text{ s}$  are ruled out for  $M_A \sim 17 - 18 \text{ MeV}$ . This has already been translated into constraints on a reduced pseudoscalar-fermion Yukawa coupling  $C_{\text{Aff}}$  in [13], where  $C_{\text{Aff}} = \xi_e$  in our notation. Following [13],  $0.4 \lesssim C_{\text{Aff}} \lesssim 4$  is ruled out by this constraint. Since  $\xi_e < 0.4$  is incompatible with (3.16), one is left with

$$\xi_e \stackrel{!}{\gtrsim} 4. \quad (4.2)$$

This constraint leads automatically to the satisfaction of the lower bound  $\xi_e \gtrsim 3.5$  from  $B^0 \rightarrow K_S^0 + \text{invisible}$ , as well as to a short enough decay length (3.16) for the Atomki pair spectrometer experiment.

Another potentially relevant experiment is the proton beam dump on copper CHARM experiment [32]. In [32] constraints were derived assuming that the production cross section and decay length of light pseudoscalars correspond to the one of axions, which is not the case here. Relevant is the analysis in [11] which uses the production of light pseudoscalars in  $K \rightarrow \pi + A$  and  $B \rightarrow X + A$  decays. For universally rescaled Yukawa couplings the region  $g_Y \gtrsim 1.5$  or  $\xi \gtrsim 1$  satisfies the constraints, since then the decay length of  $A$  is too short to reach the decay region of the CHARM experiment. This constraint does not supersede the one in (4.2).

The electron beam dump experiment E137 at SLAC [33] was analysed in terms of a decay constant  $F$  of leptophilic pseudo-Nambu-Goldstone bosons in [14]. From [14] one finds that  $F \lesssim 100 \text{ GeV}$  is allowed which corresponds, with  $\frac{1}{F} = \frac{\xi}{v}$ , to  $\xi \gtrsim 2.5$  leading again to a short decay length. Again this constraint does not supersede the one in (4.2).

Constraints from the additional electron beam dump experiments SLAC E141 [34] and Fermilab E774 [35] do not apply for  $M_A \sim 17 \text{ MeV}$ .

Since beam dump experiments are not sensitive to short decay lengths/large couplings by construction one may ask whether there are any upper limits on  $\xi_e$ . Tree level processes mediated by  $A$  with Higgs-like Yukawa couplings (even if rescaled by  $\xi_e \gtrsim 4$ ) compete with flavour conserving electroweak processes with couplings of  $\mathcal{O}(1)$ . Compared to pure electromagnetic processes at eV scales its contributions are suppressed additionally by  $(\text{eV}/M_A)^4$ . Whereas weak upper limits on  $\xi_e$  could certainly be derived from tree level processes, it is thus not astonishing that presently discussed limits on Yukawa couplings of  $A$  [11, 13, 14] rely on loop-induced flavour changing processes (and the muon anomalous moment). However, in all these cases additional BSM particles must contribute in order to restore electroweak gauge invariance. Since these can cancel the  $A$ -contribution for any  $\xi_e$  in principle, the upper limit on  $\xi_e$  depends on the amount of finetuning one is willing to tolerate which depends, however, on the UV-complete model under consideration.

## 5 Summary and conclusions

We studied for which range of Yukawa couplings — parametrized in terms of rescaled Yukawa couplings of a SM Higgs boson — a pseudoscalar with a mass of  $\sim 17 \text{ MeV}$  can

explain the anomaly observed in the Atomki pair spectrometer experiment. The production rate relative to photon emission in  ${}^8\text{Be}^*$  decays was estimated in the nuclear shell model (neglecting, amongst others, isospin-breaking effects) leading to  $\xi_u + \xi_d \approx 0.6$ ; a larger value is likely if isospin-breaking effects as discussed in [3] are important. A decay length short enough for the Atomki pair spectrometer experiment requires  $\xi_e \gtrsim 1$ .

Such a light pseudoscalar can generate flavour changing neutral currents which are constrained notably by  $K \rightarrow \pi + X$  decays. Here cancellations among the various (model dependent) loop contributions to the  $A - s - d$  vertex, at least at the 10% level, must be assumed. The dominant constraint on  $\xi_e$  is  $\xi_e \gtrsim 4$  from the electron beam dump experiment [31].

Light pseudoscalars can appear in models with extended Higgs sectors (including singlets) in which an approximate ungauged global symmetry is spontaneously broken. Examples are two-Higgs-doublet models of type II with a singlet as the NMSSM near the Peccei-Quinn or  $R$ -symmetry limit, in which case one obtains  $\xi_d \sim \xi_e$ . On the one hand, given the quite irrevocable constraints on  $\xi_e$ , this relation could only be maintained if our result for  $\xi_u + \xi_d$  is misleading by an order of magnitude due to the neglect of isospin breaking, which is not excluded. On the other hand, larger values for  $\xi_d$  would aggravate the required tuning to suppress  $K \rightarrow \pi + A$  decays. If these conditions are satisfied, models for light pseudoscalars from extended Higgs sectors could explain the anomaly observed in the Atomki pair spectrometer experiment.

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